

KUCP-71  
HUPD-9411  
August, 1994

# Spin Structure Function $g_2(x, Q^2)$ and Twist-3 Operators in QCD

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## Abstract

We investigate the spin structure function  $g_2(x, Q^2)$  in the framework of the operator product expansion and the renormalization group. The twist-3 operators appearing in QCD are examined and their relations are studied. It is noted that operators proportional to equation of motion appear in the operator mixing through renormalization, which can be studied from the relevant Green's functions. We also note that the coefficient functions can be properly fixed after the choice of independent operators.

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\*Supported in part by the Monbusho Grant-in-Aid for Scientific Research No. C-05640351.

†Supported in part by the Monbusho Grant-in-Aid for Scientific Research No. 050076

‡Supported in part by the Monbusho Grant-in-Aid for Scientific Research No. C-06640392 and No. 06221239

Recently there has been much interest in the nucleon spin structure functions which can be measured by the polarized deep inelastic leptonproduction. The nucleon spin structure is described by the two spin structure functions  $g_1(x, Q^2)$  and  $g_2(x, Q^2)$ . The data from the EMC Collaboration [1] on  $g_1$  have prompted many authors to reanalyze the QCD effects on this process and recent data from SMC at CERN [2] and E142 at SLAC [3] have attracted much more attention in connection with the Bjorken sum rule. On the other hand, the experiment on  $g_2$  is expected to be performed at CERN, SLAC and DESY in the near future.

For  $g_1$ , the leading contribution comes only from twist-2 operators in the operator product expansion (OPE), while for  $g_2$  the twist-3 operators also contribute in the leading order of  $1/Q^2$  in addition to the twist-2 operators [4]. This leads to new features which do not appear in the analyses of other structure functions.

After the old papers [5, 6, 7] which discussed the QCD effects on the polarized process, the above problems have been addressed by many authors [8, 9]. It turns out that the procedure to obtain the QCD corrections to the Wilson's coefficient functions of each operator is not so straightforward for  $g_2$ . The presence of the operators proportional to the *equation of motion* brings about some complication for the twist-3 operators, through the course of renormalization, which is a general feature characteristic to the higher-twist operators [10].

In this paper we focus our attention on the structure function  $g_2(x, Q^2)$ , and study twist-3 operators in QCD in the framework of OPE and renormalization group. We examine the operator mixing under the renormalization and point out that the coefficient functions depend upon the choice of the basis of the independent operators.

Let us start with the basic formalism based on the operator product expansion for analyzing the polarized leptonproduction [6]. The hadronic tensor  $W_{\mu\nu}$  is given by the absorptive part of the forward virtual Compton amplitude. It's anti-symmetric part

$W_{\mu\nu}^A$  can be expressed by the two spin-dependent structure functions  $g_1$  and  $g_2$  as :

$$W_{\mu\nu}^A = \varepsilon_{\mu\nu\lambda\sigma} q^\lambda \left\{ s^\sigma \frac{1}{p \cdot q} g_1(x, Q^2) + (p \cdot q s^\sigma - q \cdot s p^\sigma) \frac{1}{(p \cdot q)^2} g_2(x, Q^2) \right\},$$

where  $q$  is the virtual photon momentum and  $p$  is the nucleon momentum.  $x$  is the Bjorken variable  $x = Q^2/2p \cdot q = Q^2/2M\nu$ ,  $p \cdot q = M\nu$  and  $q^2 = -Q^2$ .  $M$  is the nucleon mass and  $s^\mu = \bar{u}(p, s)\gamma^\mu\gamma_5 u(p, s)$  is the covariant spin vector.

Applying OPE to the product of two electromagnetic currents, we get for the anti-symmetric part,

$$i \int d^4x e^{iq \cdot x} T(J_\mu(x) J_\nu(0))^A = -i \varepsilon_{\mu\nu\lambda\sigma} q^\lambda \sum_{n=1,3,\dots} \left( \frac{2}{Q^2} \right)^n q_{\mu_1} \cdots q_{\mu_{n-1}} \times \left\{ E_q^n R_q^{\sigma\mu_1 \cdots \mu_{n-1}} + \sum_j E_j^n R_j^{\sigma\mu_1 \cdots \mu_{n-1}} \right\}, \quad (1)$$

where  $R_i^n$ 's are the composite operators and  $E_i^n$ 's the corresponding coefficient functions. In (1),  $R_q$  represents the twist-2 operators and the other operators inside the summation over  $j$  are the twist-3 operators. For simplicity, let us consider the flavor non-singlet case. (In the following expressions, we suppress the flavor matrices  $\lambda_i$  for the quark field  $\psi$ ) The explicit forms of the twist-2 operators are given by

$$R_q^{\sigma\mu_1 \cdots \mu_{n-1}} = i^{n-1} \bar{\psi} \gamma_5 \gamma^{\{\sigma} D^{\mu_1} \cdots D^{\mu_{n-1}\}} \psi - (\text{traces}),$$

where  $\{ \}$  means symmetrization over the Lorentz indices and  $-(\text{traces})$  stands for the subtraction of the trace terms to make the operators traceless, which will be suppressed in the following. For the twist-3 operators, we have

$$R_F^{\sigma\mu_1 \cdots \mu_{n-1}} = \frac{i^{n-1}}{n} \left[ (n-1) \bar{\psi} \gamma_5 \gamma^\sigma D^{\{\mu_1} \cdots D^{\mu_{n-1}\}} \psi - \sum_{l=1}^{n-1} \bar{\psi} \gamma_5 \gamma^{\mu_l} D^{\{\sigma} D^{\mu_1} \cdots D^{\mu_{l-1}} D^{\mu_{l+1}} \cdots D^{\mu_{n-1}\}} \psi \right], \quad (2)$$

$$R_m^{\sigma\mu_1 \cdots \mu_{n-1}} = i^{n-2} m \bar{\psi} \gamma_5 \gamma^\sigma D^{\{\mu_1} \cdots D^{\mu_{n-2}} \gamma^{\mu_{n-1}\}} \psi, \quad (3)$$

$$R_k^{\sigma\mu_1 \cdots \mu_{n-1}} = \frac{1}{2n} (V_k - V_{n-1-k} + U_k + U_{n-1-k}), \quad (4)$$

where  $m$  in (3) represents the quark mass (matrix). The operators in (4) contain the gluon field strength  $G_{\mu\nu}$  and the dual tensor  $\tilde{G}_{\mu\nu} = \frac{1}{2}\varepsilon_{\mu\nu\alpha\beta}G^{\alpha\beta}$  explicitly and given by

$$\begin{aligned} V_k &= i^n g S \bar{\psi} \gamma_5 D^{\mu_1} \dots G^{\sigma\mu_k} \dots D^{\mu_{n-2}} \gamma^{\mu_{n-1}} \psi, \\ U_k &= i^{n-3} g S \bar{\psi} D^{\mu_1} \dots \tilde{G}^{\sigma\mu_k} \dots D^{\mu_{n-2}} \gamma^{\mu_{n-1}} \psi, \end{aligned}$$

where  $S$  means the symmetrization over  $\mu_i$  and  $g$  is the QCD coupling constant.

It is by now well-known [8, 9] that the operators (2 - 4) are related through the operators which are proportional to the *equation of motion* ( EOM operators ),

$$\begin{aligned} R_{eq}^{\sigma\mu_1\dots\mu_{n-1}} &= i^{n-2} \frac{n-1}{2n} S [\bar{\psi} \gamma_5 \gamma^\sigma D^{\mu_1} \dots D^{\mu_{n-2}} \gamma^{\mu_{n-1}} (i \not{D} - m) \psi \\ &\quad + \bar{\psi} (i \not{D} - m) \gamma_5 \gamma^\sigma D^{\mu_1} \dots D^{\mu_{n-2}} \gamma^{\mu_{n-1}} \psi]. \end{aligned}$$

It is not difficult to obtain the following relation :

$$R_F^{\sigma\mu_1\dots\mu_{n-1}} = \frac{n-1}{n} R_m^{\sigma\mu_1\dots\mu_{n-1}} + \sum_{k=1}^{n-2} (n-1-k) R_k^{\sigma\mu_1\dots\mu_{n-1}} + R_{eq}^{\sigma\mu_1\dots\mu_{n-1}}. \quad (5)$$

This equation means that these operators are not all independent but constrained by this relation, and therefore the renormalization procedure for the twist-3 operators becomes rather complicated. Although we realize that physical matrix elements of the EOM operators vanish [10], we keep them to study the operator mixing through renormalization because their off-shell Green's function do not vanish and are relevant for our analyses shown below.

The anomalous dimensions which enter into the renormalization group equation for the coefficient functions, are obtained from the renormalization constant for the composite operators. The anomalous dimensions of twist-3 operators in QCD at the one-loop level have been studied in ref.[8]. Here we reanalyze the operator mixing problems by keeping the EOM operators. We examine what happens to the renormalization if there are several operators which are related by constraints. Let us first

illustrate the points with the scalar field theory for simplicity. Assume that we have two composite operators  $R_1$  and  $R_2$  which are related through the equation of motion:

$$R_1 = R_2 + E, \quad (6)$$

where  $E$  is a EOM operator. The general form of  $E$  will be  $E = A(\phi) \delta S / \delta \phi$  where  $S$  is the action and  $A(\phi)$  is a some function of  $\phi$ . Since we have three composite operators with a constraint (6), we can choose any two operators among three as an independent basis.

First we choose  $R_1$  and  $R_2$ . Then renormalization takes the following form:

$$\begin{pmatrix} R_1 \\ R_2 \end{pmatrix}_R = \begin{pmatrix} Z_{11} & Z_{12} \\ Z_{21} & Z_{22} \end{pmatrix} \begin{pmatrix} R_1 \\ R_2 \end{pmatrix}_B. \quad (7)$$

Here the suffix  $R$  ( $B$ ) denotes the renormalized (bare) operator. Next, we change the operator basis to  $R_i$  ( $i = 1$  or  $2$ ) and  $E$ . Then the renormalization matrix becomes triangular [11],

$$\begin{pmatrix} R_i \\ E \end{pmatrix}_R = \begin{pmatrix} Z_{ii}^E & Z_{iE} \\ 0 & Z_{EE} \end{pmatrix} \begin{pmatrix} R_i \\ R_E \end{pmatrix}_B, \quad (8)$$

since the counterterm to the operator  $E$  should vanish by the equation of motion. From (7) and (8), we can derive the following relations between the renormalization constants  $Z$ 's:

$$\begin{aligned} Z_{11}^E = Z_{11} + Z_{12} &= Z_{21} + Z_{22} = Z_{22}^E, \\ Z_{EE} = Z_{22} - Z_{12} = Z_{11} - Z_{21}, &\quad Z_{1E} = -Z_{12}, \quad Z_{2E} = Z_{21}. \end{aligned} \quad (9)$$

Eq.(9) is a consistency condition for the renormalization of the operators satisfying (6).

Here we note that the arbitrariness in the choice of the operator basis does not enter into the physical quantities. What we really need is the physical matrix element of operators. This fact can be explicitly confirmed by noting that the physical matrix element of the operator  $E$  vanishes [10, 11]:  $\langle p|E|p' \rangle = 0$ . Taking the physical matrix

element of (8), we conclude that the renormalization constant  $Z_{ii}^E$  is only relevant. We reach the same conclusion also from (7) by taking account of the fact:  $\langle p|R_1|p'\rangle = \langle p|R_2|p'\rangle$  and (9). The explicit one loop calculation shows that for the simplest twist-3 operator  $\bar{\psi}D_\mu D_\nu\psi$ , as discussed by Jaffe [9], the above argument actually holds [12].

Now we go back to the case of polarized deep inelastic scattering and give explicit results at the one-loop level for  $n = 3$  as an example. In this case we have four operators with the constraint (5).

$$R_F = \frac{2}{3}R_m + R_1 + R_{eq}, \quad (10)$$

where Lorentz indices of operators are suppressed. Now we choose  $R_F, R_m$ , and  $R_1$  as independent operators and eliminate EOM operator  $R_{eq}$ . In order to renormalize above operators, it turns out that a gauge non-invariant EOM operator  $R_{eq1}$  comes into play,

$$R_{eq1}^{\sigma\mu_1\mu_2} = i\frac{1}{3}S[\bar{\psi}\gamma_5\gamma^\sigma\partial^{\mu_1}\gamma^{\mu_2}(i\not{D} - m)\psi + \bar{\psi}(i\not{D} - m)\gamma_5\gamma^\sigma\partial^{\mu_1}\gamma^{\mu_2}\psi].$$

Its presence is allowed because it vanishes by the equation of motion [11]. We found the following renormalization constant for the composite operators,

$$\begin{pmatrix} R_F \\ R_1 \\ R_m \\ R_{eq1} \end{pmatrix}_R = \begin{pmatrix} Z_{11} & Z_{12} & Z_{13} & Z_{14} \\ Z_{21} & Z_{22} & Z_{23} & Z_{24} \\ 0 & 0 & Z_{33} & 0 \\ 0 & 0 & 0 & Z_{44} \end{pmatrix} \begin{pmatrix} R_F \\ R_1 \\ R_m \\ R_{eq1} \end{pmatrix}_B \quad (11)$$

where  $Z_{ij}$  are given in the dimensional regularization  $D = 4 - 2\varepsilon$ :

$$Z_{ij} \equiv \delta_{ij} + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} z_{ij}$$

with

$$\begin{aligned} z_{11} &= \frac{7}{6}C_2(R) + \frac{3}{8}C_2(G), & z_{12} &= -\frac{3}{2}C_2(R) + \frac{21}{8}C_2(G), \\ z_{13} &= 3C_2(R) - \frac{1}{4}C_2(G), & z_{14} &= -\frac{3}{8}C_2(G), \\ z_{21} &= \frac{1}{6}C_2(R) - \frac{1}{8}C_2(G), & z_{22} &= -\frac{1}{2}C_2(R) + \frac{25}{8}C_2(G), \\ z_{23} &= -\frac{1}{3}C_2(R) + \frac{1}{12}C_2(G), & z_{24} &= \frac{1}{8}C_2(G), \\ z_{33} &= 6C_2(R), & z_{44} &= 0. \end{aligned} \quad (12)$$

The quadratic Casimir operators are  $C_2(R) = 4/3$  and  $C_2(G) = 3$  for the case of QCD. In the above calculations, it is necessary to compute the *off-shell Green's func-*

tions of the composite operators. Otherwise some informations on the renormalization constants associated with the EOM operator will be lost.

As explained before, one can choose other operators as an independent basis using (10). The consistency conditions corresponding to (9) read in this case,

$$\begin{aligned} z_{11} + z_{12} &= z_{21} + z_{22}, & \frac{2}{3}z_{11} + z_{13} &= \frac{2}{3}z_{21} + z_{23} + \frac{2}{3}z_{33}, \\ z_{13} - \frac{2}{3}z_{12} &= z_{23} - \frac{2}{3}z_{22} + \frac{2}{3}z_{33}. \end{aligned}$$

These equalities are indeed satisfied with (12).

The authors in ref.[8] have discarded the fermion bilinear operator  $R_F^{\sigma\mu_1\cdots\mu_{n-1}}$  in their analyses. For the present case of  $n = 3$ , the renormalization constant matrix becomes

$$\begin{pmatrix} R_1 \\ R_m \\ R_{eq} \\ R_{eq1} \end{pmatrix}_R = \begin{pmatrix} Z_{21} + Z_{22} & \frac{2}{3}Z_{21} + Z_{23} & Z_{21} & Z_{24} \\ 0 & Z_{33} & 0 & 0 \\ 0 & 0 & Z_{11} - Z_{21} & Z_{14} - Z_{24} \\ 0 & 0 & 0 & Z_{44} \end{pmatrix} \begin{pmatrix} R_1 \\ R_m \\ R_{eq} \\ R_{eq1} \end{pmatrix}_B, \quad (13)$$

where  $Z_{ij}$  are defined in (11). In this basis, our results agree with those in ref.[8]. Although this choice of basis might be economical especially for general  $n$ , it is never compulsory. One can eliminate any operator as well. Any choice of basis leads to a unique prediction for the moment with the properly determined coefficient functions discussed below.

To make a prediction for the moments, we must determine the coefficient functions corresponding to the “hard part” of the process. The coefficient functions at the tree level can be obtained by considering the short distance expansion of the current product in the presence of external gauge field  $A_\mu^a$ . Up to twist-3 operators, the OPE reads [8, 9]

$$\begin{aligned} & i \int d^4x e^{iq \cdot x} T(J_\mu(x) J_\nu(0))^A \\ &= -i \varepsilon_{\mu\nu\lambda\sigma} q^\lambda \sum_{n=1,3,\dots} \left( \frac{2}{Q^2} \right)^n q_{\mu_1} \cdots q_{\mu_{n-1}} \{ R_q^{\sigma\mu_1\cdots\mu_{n-1}} + R_F^{\sigma\mu_1\cdots\mu_{n-1}} \}. \end{aligned} \quad (14)$$

Here if we use (5),  $R_F^{\sigma\mu_1\cdots\mu_{n-1}}$  can be eliminated in terms of other operators. From this observation, it is inferred that the (tree-level) coefficient functions depend upon

the choice of the independent operators. In the basis of independent operators which includes  $R_F^n$ , we conclude that at the tree-level,

$$E_q^n(\text{tree}) = E_F^n(\text{tree}) = 1, \quad E_m^n(\text{tree}) = E_j^n(\text{tree}) = 0. \quad (15)$$

On the other hand, if we eliminate  $R_F^n$  we have,

$$E_q^n(\text{tree}) = 1, \quad E_m^n(\text{tree}) = \frac{n-1}{n}, \quad E_j^n(\text{tree}) = n-1-j. \quad (16)$$

Now we shall write down the moment sum rule for  $g_2$ . Define the matrix elements of composite operators between nucleon states with momentum  $p$  and spin  $s$  by

$$\langle p, s | R_q^{\sigma\mu_1 \dots \mu_{n-1}} | p, s \rangle = -a_n s^{\{\sigma} p^{\mu_1} \dots p^{\mu_{n-1}\}} \quad (17)$$

$$\langle p, s | R_F^{\sigma\mu_1 \dots \mu_{n-1}} | p, s \rangle = -\frac{n-1}{n} d_n (s^\sigma p^{\mu_1} - s^{\mu_1} p^\sigma) p^{\mu_2} \dots p^{\mu_{n-1}} \quad (18)$$

$$\langle p, s | R_m^{\sigma\mu_1 \dots \mu_{n-1}} | p, s \rangle = -e_n (s^\sigma p^{\mu_1} - s^{\mu_1} p^\sigma) p^{\mu_2} \dots p^{\mu_{n-1}} \quad (19)$$

$$\langle p, s | R_k^{\sigma\mu_1 \dots \mu_{n-1}} | p, s \rangle = -f_n^k (s^\sigma p^{\mu_1} - s^{\mu_1} p^\sigma) p^{\mu_2} \dots p^{\mu_{n-1}} \quad (20)$$

$$\langle p, s | R_{eq}^{\sigma\mu_1 \dots \mu_{n-1}} | p, s \rangle = 0. \quad (21)$$

Our normalization in the above definition is such that for free quark target  $a_n = d_n = e_n = 1$ . On the other hand,  $f_n^k = \mathcal{O}(g^2)$ . Using (17 - 21), we can write down the moment sum rule for  $g_2$ .

$$M_n \equiv \int_0^1 dx x^{n-1} g_2(x, Q^2) = -\frac{n-1}{2n} [a_n E_q^n(Q^2) - d_n E_F^n(Q^2)] + \frac{1}{2} [e_n E_m^n(Q^2) + \sum_{j=1}^{n-2} f_n^j E_j^n(Q^2)]. \quad (22)$$

with the following constraint from (5),

$$\frac{n-1}{n} d_n = \frac{n-1}{n} e_n + \sum_{j=1}^{n-2} (n-1-j) f_n^j.$$

It should be stressed that the explicit form of the  $Q^2$  evolution of each term in (22) depends on the choice of operator basis. The anomalous dimensions as well as the



coefficient functions  $E_j^n$  take different forms depending on the basis. However the moment itself remain the same as one expects. For the case of  $n = 3$ , we have explicitly checked that any choice of the basis leads to the same expression for the moment.

Here we note that the coefficient functions for the EOM operators do not mix with those for the other operators. This is due to the triangular structure of the anomalous dimension matrix in their renormalization-group equation. This fact is essential to the equivalence of the moments for any basis.

Let us now consider the moments of  $g_2$  in the basis of  $R_j$  and  $R_m$ . In this case,

$$M_n = -\frac{n-1}{2n}a_n E_q^n(Q^2) + \frac{1}{2}\left[e_n E_m^n(Q^2) + \sum_{j=1}^{n-2} f_n^j E_j^n(Q^2)\right]. \quad (23)$$

with the coefficient functions given in (16). We shall show below that (23) indeed holds in the leading order of  $\ln Q^2$ .

The operator mixing of gluon-field dependent operators  $R_j$  and the mass dependent operator  $R_m$  through the renormalization reads neglecting the EOM operators,

$$\begin{pmatrix} R_j \\ R_m \end{pmatrix}_R = \begin{pmatrix} Z_{ji} & Z_{jm} \\ 0 & Z_{mm} \end{pmatrix} \begin{pmatrix} R_i \\ R_m \end{pmatrix}_B.$$

By evaluating the Green's functions of these operators with incoming and outgoing off-shell quark states, we obtain the off-diagonal element of the renormalization constant matrix  $Z_{jm}$  from which the anomalous dimension reads,

$$\gamma_{mj} \equiv -\frac{g^2}{16\pi^2} \frac{8C_2(R)}{n} \frac{1}{j(j+1)(j+2)} \equiv \frac{g^2}{16\pi^2} \gamma_{mj}^0. \quad (24)$$

This result is in disagreement with the one given in the fifth reference in [8].

Now note that the Compton scattering amplitude off the on-shell massive quark target has been calculated [6], in the leading order of  $\ln Q^2$ , to be

$$M_n = \frac{1}{2} \frac{g^2}{16\pi^2} C_2(R) \left(-2 + \frac{4}{n+1}\right) \ln Q^2 + \dots \quad (25)$$

Whereas, using the perturbative solution for the renormalization-group equation, the right-hand side of (23) becomes for the quark matrix elements,

$$\text{RHS of (23)} = -\frac{1}{2} \frac{g^2}{16\pi^2} \left( \frac{n-1}{n} \left( -\frac{1}{2} \right) \gamma_q^0 E_q^n(\text{tree}) + \frac{1}{2} \gamma_{mj}^0 E_j(\text{tree}) \right) \ln Q^2 + \dots \quad (26)$$

Putting  $\gamma_{mj}^0$  of (24) and the following expression for the anomalous dimensions,

$$\gamma_q^0 = 2C_2(R) \left[ 1 - \frac{2}{n(n+1)} + 4 \sum_{j=2}^n \frac{1}{j} \right], \quad \gamma_{mm}^0 = 8C_2(R) \sum_{j=1}^{n-1} \frac{1}{j}.$$

into the above equation with the tree level coefficient functions (16), we find (26) coincides with (25). Thus we confirm that our results (24,16) are consistent with the moment sum rule.

Finally let us comment on the Burkhardt-Cottingham sum rule which corresponds to the first moment of  $g_2$ . We see from (2-4) that twist-3 operators can *not* be defined for  $n = 1$ . Therefore the OPE analysis suggests that the Burkhardt-Cottingham sum rule:

$$\int_0^1 dx \, g_2(x, Q^2) = 0$$

does not receive any corrections in QCD [13, 14].

To summarize, we have examined the twist-3 operators in QCD which contribute to  $g_2(x, Q^2)$ . In the renormalization, there appear operators proportional to the equation of motion, which can be studied by computing off-shell Green's functions. Because of the relationship among the twist-3 operators, we have to choose a basis of independent operators, by which the coefficient functions are properly fixed. We have also noted that the Burkhardt-Cottingham sum rule holds in the higher-order of QCD [14].

We expect that future measurements on  $g_2$  will clarify the effect of twist-3 operators which share some common feature with more general higher-twist effects in QCD.

We would like to thank Shoji Hashimoto and Ken Sasaki for discussions.

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